

Optically induced Aharonov-Bohm effect in mesoscopic rings

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We show theoretically that strong electron coupling to circularly polarized photons in non-singly-connected nanostructures results in the appearance of an artificial gauge field that changes the electron phase. The effect arises from the breaking of time-reversal symmetry and is analogous to the well-known Aharonov-Bohm phase effect. It can manifest itself in the oscillations of conductance as a function of the intensity and frequency of the illumination. The theory of the effect is elaborated for mesoscopic rings in both ballistic and diffusive regimes.

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I. INTRODUCTION

Progress in modern nanotechnologies has resulted in rapid developments in the fabrication of mesoscopic objects, including non-singly-connected nanostructures such as quantum rings. The fundamental interest attracted by these systems is caused by a wide variety of purely quantum-mechanical topological effects which can be observed in ring-like mesoscopic structures. The most notable phenomenon amongst them is the Aharonov-Bohm (AB) effect arisen from the direct influence of a vector potential on the electron phase.^{1,2} In the ballistic regime, this effect results in magnetic-flux-dependent oscillations of the conductance in ring-like structures with a period equal to the fundamental magnetic flux quantum $\Phi_0 = h/e$.³⁻⁶ In the diffusive regime, a second type of conductance oscillations with the period $\Phi_0/2$ can be observed. They are known as the Altshuler-Aronov-Spivak (AAS) oscillations and are associated with the weak localization of electrons.⁷⁻¹⁰

From a fundamental viewpoint, the AB-AAS oscillations arise from the broken time-reversal symmetry in the electron system (conducting mesoscopic ring) subjected to a magnetic flux through the ring. Namely, the flux breaks the equivalence of clockwise and counterclockwise electron rotations inside the ring, which results in the flux-controlled interference of electron waves corresponding to these rotations. The similar broken equivalence of electron motion for mutually opposite directions caused by a magnetic field can take place in various nanostructures, including quantum wells,¹¹ carbon nanotubes¹² and hybrid semiconductor/ferromagnet nanostructures.¹³ However, the time-reversal symmetry can be broken not only by a magnetic flux but also by a circularly polarized electromagnetic field. Indeed, the field breaks the symmetry since time reversal turns clockwise polarized photons into counterclockwise polarized ones and vice versa. As a result, the electron coupling to circularly polarized photons can change electron energy spectrum of quantum rings.¹⁴ Therefore, phenomena similar to the AB effect can take place in ring-like electronic

systems interacting with a circularly polarized electromagnetic field. We will show below that the conductance of these electron-photon systems can exhibit oscillations which are formally equivalent to the AB-AAS oscillations induced by a magnetic flux. The phenomenon can be described in terms of an artificial $U(1)$ gauge field generated by the strong coupling between electron and circularly polarized photons. The theory of such optically-induced AB effect, which lies at the border between condensed-matter physics and quantum optics, is developed in this paper.

The paper is organized as follows. In the Section II, we introduce the Schrödinger problem describing the electron interaction with circularly polarized photons in mesoscopic rings. The Section III is devoted to derivation of an artificial $U(1)$ gauge field arisen from the strong electron-photon coupling in the rings. In the Section IV, AB-AAS oscillations of conductance caused by the gauge field are analyzed.

II. THE MODEL

Let us consider the conventional model of an electron interference device (see, e.g., Refs.¹⁵⁻¹⁷) consisting of an one-dimensional mesoscopic ring with radius R and two one-dimensional leads which are connected at the quantum point contacts (see Fig. 1a). Generally, the phase shift between the clockwise and counterclockwise traveling electron waves,

$$\Delta\phi = \phi_+ - \phi_-, \quad (1)$$

can be nonzero: The shift can be caused by the application of an external magnetic field (AB effect) or result from spin-orbit interaction.¹⁸⁻²¹ Experimentally, it can be detected by measuring the field-dependent oscillations of the conductance of the device.

In order to write the phase shift (1) as a function of the field parameters, we have to consider the electron energy spectrum of an isolated ring subjected to an electromagnetic field with the vector potential \mathbf{A} . If the field

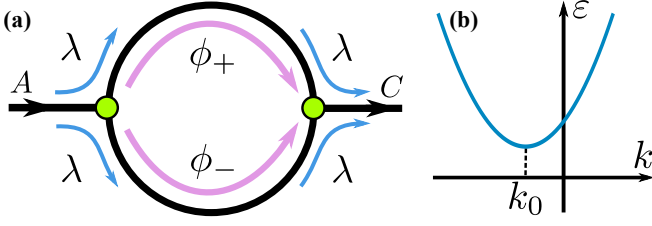


FIG. 1: (Color online) **(a)** The scheme of the electron interference device consisting of an one-dimensional mesoscopic ring which is connected with two one-dimensional leads at the quantum point contacts (QPCs). An electron wave, which enters into the device with the amplitude $A = 1$, is split between the two different paths with the transmission amplitudes λ and exits the device with the amplitude C . The phase shift of the electron waves traveling clockwise and counterclockwise inside the ring, $\Delta\phi = \phi_+ - \phi_-$, arises from the coupling to an external electromagnetic field; **(b)** The scheme of the electron energy spectrum $\varepsilon(k)$ in a mesoscopic ring subjected to a circularly polarized electromagnetic field or a stationary magnetic field. The spectrum is shifted along the k axis by the wave vector k_0 which depends on the parameters of the field.

is time-independent, then the electron energy spectrum can be found from the stationary Schrödinger equation with the Hamiltonian

$$\hat{H}_0 = \frac{1}{2m_e} (\hat{p}_\varphi - eA_\varphi)^2, \quad (2)$$

where φ is the electron angular coordinate in the ring, $\hat{p}_\varphi = -i(\hbar/R)\partial/\partial\varphi$ is the operator of electron momentum in the ring, e is the electron charge, and m_e is the effective electron mass in the ring. Particularly, in the well-known case of a stationary magnetic field, \mathbf{B} , directed perpendicularly to the ring plane, the electron energy spectrum of the ring has the form

$$\varepsilon(m) = \frac{\hbar^2}{2m_e R^2} \left(m + \frac{\Phi}{\Phi_0} \right)^2,$$

where $m = 0, \pm 1, \pm 2, \pm 3, \dots$ is the electron angular momentum along the ring axis, and $\Phi = B\pi R^2$ is the magnetic flux through the ring. In the considered case of a mesoscopic ring, it is convenient to rewrite this spectrum as

$$\varepsilon(k) = \frac{\hbar^2}{2m_e} \left(k + \frac{\Phi}{R\Phi_0} \right)^2, \quad (3)$$

where $k = m/R$ is the electron wave vector along the ring. Graphically, the energy spectrum (3) can be pictured as a parabola shifted along the k axis by the wave vector

$$k_0 = -\Phi/R\Phi_0 \quad (4)$$

(see Fig. 1b). Formally, just the wave vector (4) defines the nonzero phase shift (1) since $\Delta\phi = 2\pi Rk_0$.

Any electromagnetic field, which results to such a shifted electron energy spectrum with $k_0 \neq 0$, can generate the oscillations of conductance of the considered electron interference device. However, in the case of a time-dependent electromagnetic field with the vector potential $A_\varphi(t)$, the Schrödinger equation with the Hamiltonian (2) is non-stationary and cannot be used to find the electron energy spectrum. The regular approach to solve this quantum-mechanical problem should be based on the conventional methodology of quantum optics.^{25,26} Namely, we have to consider the system “electrons in the ring + electromagnetic field” as a whole and to write the Hamiltonian of this electron-photon system. If the field frequency lies far from the resonant frequencies of the electron subsystem (i.e. the field is purely “dressing”), then the energy spectrum of the electron-photon system can be written as a sum of field energy and energy of the electrons strongly coupled to the field (dressed electrons). This energy spectrum of dressed electrons will be responsible for all electron characteristics of the ring subjected to the strong high-frequency electromagnetic field.

The Hamiltonian (2) is written as a function on the vector potential $A_\varphi(t)$ which depends on the gauge. In order to rewrite the Hamiltonian in gauge invariant form, let us apply the unitary transformation²⁵

$$U = \exp \left(\frac{ieR}{\hbar} \int A_\varphi(t) d\varphi \right),$$

where the indefinite integral over the angle φ should be treated as an anti-derivative of the integrand. Then the transformed electron Hamiltonian (2),

$$\hat{H}'_0 = U^\dagger \hat{H}_0 U - i\hbar U^\dagger \frac{\partial U}{\partial t},$$

takes the form

$$\hat{H}'_0(E_\varphi) = \frac{\hat{p}_\varphi^2}{2m_e} - eR \int E_\varphi d\varphi, \quad (5)$$

where $E_\varphi = -\partial A_\varphi(t)/\partial t$ is the angular component of the electric field which does not depend on the field gauge. Although the interaction of electrons in ring-like structures with an electric field was considered previously (see, e.g., Ref. [22,23]), phase-shift phenomena caused by a high-frequency field have so far escaped attention. Considering the problem within the conventional quantum-field approach,²⁴⁻²⁶ the classical electric field in the Hamiltonian (5), \mathbf{E} , should be replaced with the field operator, $\hat{\mathbf{E}}$. Then the complete electron-photon Hamiltonian reads

$$\hat{H} = \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} \hat{a}_{\mathbf{q}}^\dagger \hat{a}_{\mathbf{q}} + \hat{H}'_0(\hat{E}_\varphi), \quad (6)$$

where the first term describes the field energy, \mathbf{q} is the photon wave vector, $\omega_{\mathbf{q}}$ is the photon frequency, $\hat{a}_{\mathbf{q}}^\dagger$ and $\hat{a}_{\mathbf{q}}$ are the photon operators of creation and annihilation respectively, and the summation is assumed to be

performed over all photon modes of the electromagnetic field. If the ring is subjected to a monochromatic circularly polarized electromagnetic wave propagating perpendicularly to the ring, the Hamiltonian (6) takes the form

$$\hat{\mathcal{H}} = \hbar\omega\hat{a}^\dagger\hat{a} + \frac{\hat{p}_\varphi^2}{2m_e} - ieR\sqrt{\frac{\hbar\omega}{4\epsilon_0 V}}(e^{i\varphi}\hat{a} - e^{-i\varphi}\hat{a}^\dagger), \quad (7)$$

where ω is the field frequency. Considering the last term in the Hamiltonian (7) as a perturbation, we can apply the approach developed in Ref. 14 to solve the electron-photon Schrödinger equation with this Hamiltonian. From experimental viewpoint, the most relevant case corresponds to the ring exposed to a classically strong laser-generated electromagnetic field. Just such a strong electromagnetic field will be under consideration in the following. In contrast to the case of a ring interacting with a weak photon mode inside a cavity,^{27,28} an amplitude of the strong field does not depend on the electron-photon interaction. As a result, the energy spectrum of dressed electrons in the ring can be found as an expansion in terms of the dressing field amplitude E_0 . Assuming the inequality $|e|E_0/m_e R\omega^2 \ll 1$ to be satisfied and accounting for terms squared in the field amplitude only, the energy spectrum of dressed electrons in the ring can be written as

$$\varepsilon(k) = \frac{\hbar^2 k^2}{2m_e} + \frac{\hbar e^2 E_0^2}{2m_e^2 R\omega^3} k. \quad (8)$$

It should be noted that the Hamiltonian (7) describes electrons in an isolated ring, where the electron lifetime is $\tau \rightarrow \infty$. In the interference device pictured in Fig. 1a, this lifetime is the traveling time of an electron from one QPC to the other one, i.e. $\tau \sim \pi R/v_F$, where v_F is the Fermi velocity of an electron in the ring. Therefore, the developed theory is consistent if the field frequency, ω , is large enough to satisfy the condition $2\pi/\omega\tau \ll 1$ which allows one to consider the incident electromagnetic field as a dressing field.

The energy spectrum (8) has the form plotted in Fig. 1b with

$$k_0 = -\frac{e^2 E_0^2}{2m_e \hbar R\omega^3}. \quad (9)$$

It follows from the comparison of Eqs. (4) and (9) that the high-frequency electromagnetic field results in the same phase shift (1) as an effective magnetic flux

$$\Phi_{\text{eff}} = \frac{|e|\pi E_0^2}{m_e \omega^3}. \quad (10)$$

Let us show that the effective magnetic flux (10) can be described in terms of an artificial $U(1)$ gauge field with the vector potential

$$A_\varphi^{\text{eff}} = \frac{\Phi_{\text{eff}}}{2\pi R} = \frac{|e|E_0^2}{2Rm_e\omega^3}, \quad (11)$$

which is produced by the strong electron-photon coupling.

III. THE ARTIFICIAL GAUGE FIELD

To describe the electron-photon coupling in the considered system, let us use the joined electron-photon space $|m, N\rangle = |\psi_m(\varphi)\rangle \otimes |N\rangle$. This corresponds to the electromagnetic field being in the state with the photon occupation number $N = 1, 2, 3, \dots$, and the electron being in the state with the wave function

$$\psi_m(\varphi) = \sqrt{1/2\pi} e^{im\varphi}, \quad (12)$$

where $m = 0, \pm 1, \pm 2, \dots$ is the electron angular momentum along the ring axis. The electron-photon states $|m, N\rangle$ are true eigenstates of the Hamiltonian of the noninteracting electron-photon system,

$$\hat{\mathcal{H}}^{(0)} = \hbar\omega\hat{a}^\dagger\hat{a} + \frac{\hat{p}_\varphi^2}{2m_e}, \quad (13)$$

and their energy spectrum is

$$\varepsilon_{m,N}^{(0)} = N\hbar\omega + \frac{\hbar^2 m^2}{2m_e R^2}. \quad (14)$$

Considering the last term in the Hamiltonian (7) as a perturbation with the matrix elements

$$\begin{aligned} \langle m', N' | \hat{U} | m, N \rangle &= -ieR\sqrt{\frac{\hbar\omega}{4\epsilon_0 V}} \left[\sqrt{N}\delta_{m,m'-1}\delta_{N,N'+1} \right. \\ &\quad \left. - \sqrt{N+1}\delta_{m,m'+1}\delta_{N,N'-1} \right] \end{aligned} \quad (15)$$

and performing trivial calculations within the first order of the perturbation theory, we can write eigenstates of the Hamiltonian (7) as

$$\begin{aligned} |\Psi_{m,N}\rangle &= \frac{\langle m+1, N-1 | \hat{U} | m, N \rangle}{\varepsilon_{m,N}^{(0)} - \varepsilon_{m+1,N-1}^{(0)}} |m+1, N-1\rangle \\ &+ \frac{\langle m-1, N+1 | \hat{U} | m, N \rangle}{\varepsilon_{m,N}^{(0)} - \varepsilon_{m-1,N+1}^{(0)}} |m-1, N+1\rangle + |m, N\rangle. \end{aligned} \quad (16)$$

Substituting Eqs. (14)–(15) into Eq. (16) and assuming the electromagnetic field to be strong ($N \gg 1$), we arrive at the expression

$$\begin{aligned} |\Psi_{m,N}\rangle &= |m, N\rangle - \frac{ieRE_0}{2} \left[\frac{|m+1, N-1\rangle}{\hbar\omega - \varepsilon_R(1+2m)} \right. \\ &\quad \left. + \frac{|m-1, N+1\rangle}{\hbar\omega + \varepsilon_R(1-2m)} \right], \end{aligned} \quad (17)$$

where $E_0 = \sqrt{N\hbar\omega/\epsilon_0 V}$ is the classical amplitude of electric field, and $\varepsilon_R = \hbar^2/2m_e R^2$ is the characteristic electron energy in the ring. Taking into account Eq. (12), we can rewrite the basis electron-photon states as $|m \pm 1, N\rangle = e^{\pm i\varphi} |m, N\rangle$. Then the eigenstates (17) takes the form

$$\begin{aligned} |\Psi_{m,N}\rangle &= |m, N\rangle - \frac{ieRE_0}{2} \left[\frac{e^{i\varphi} |m, N-1\rangle}{\hbar\omega - \varepsilon_R(1+2m)} \right. \\ &\quad \left. + \frac{e^{-i\varphi} |m, N+1\rangle}{\hbar\omega + \varepsilon_R(1-2m)} \right]. \end{aligned} \quad (18)$$

In the basis of the three electron-photon states,

$$\begin{pmatrix} |m, N+1\rangle \\ |m, N\rangle \\ |m, N-1\rangle \end{pmatrix} \quad (19)$$

the eigenstate (18) can be written formally as a vector

$$|\chi\rangle = \begin{pmatrix} -\frac{ieRE_0/2}{\hbar\omega + \varepsilon_R(1-2m)}e^{-i\varphi} \\ 1 \\ -\frac{ieRE_0/2}{\hbar\omega - \varepsilon_R(1+2m)}e^{i\varphi} \end{pmatrix}. \quad (20)$$

It should be noted that each of the basis states (19) corresponds to the same electron angular momentum m . Therefore, the influence of the electromagnetic field on the electron results only in the phase incursion describing by the exponential factors $e^{\pm i\varphi}$ in the state vector (20). Following the conventional theory of artificial gauge fields (see, e.g., Ref. 29), we can introduce the $U(1)$ field with the vector potential, $\mathbf{A}^{\text{eff}} = (i\hbar/e)\langle\chi|\nabla|\chi\rangle$, which corresponds to this phase incursion. In the case of the ring, this vector potential has the form $\mathbf{A}^{\text{eff}} = (0, 0, A_\varphi^{\text{eff}})$, where

$$A_\varphi^{\text{eff}} = \frac{i\hbar}{eR} \left\langle \chi \left| \frac{\partial}{\partial \varphi} \right| \chi \right\rangle. \quad (21)$$

Substituting Eq. (20) into Eq. (21), we arrive at the expression

$$A_\varphi^{\text{eff}} = \frac{\hbar e R E_0^2}{4} \left[\frac{1}{[\hbar\omega + \varepsilon_R(1-2m)]^2} - \frac{1}{[\hbar\omega - \varepsilon_R(1+2m)]^2} \right]. \quad (22)$$

Under the condition $\hbar\omega \gg \varepsilon_R$, the artificial vector potential (22) takes the form (11).

IV. DISCUSSION AND CONCLUSIONS

Replacing the magnetic flux Φ with the pseudo-flux (10) in known expressions which describe the oscillations of the conductance of the considered interference device, we can calculate them as follows.

First of all, let us consider the ballistic regime. In this case, the conductance is described by the Landauer formula $G = (2e^2/h)|C|^2$, where the transmission amplitude of the interference device, C , depends on the coupling between the leads and the ring. Generally, this coupling can be described by lead-to-ring and ring-to-lead transmission amplitudes, λ .^{15–17} If the reflection from one lead to itself is absent (i.e., there is no electron backscattering from QPCs), the transmission amplitude is $\lambda = \pm 1/\sqrt{2}$. This corresponds to the incoming electron wave being divided equally in the ring along the clockwise (ϕ_+) and counterclockwise (ϕ_-) paths (see Fig. 1a). In this simplest case, the replacement $\Phi \rightarrow \Phi_{\text{eff}}$ in the expression

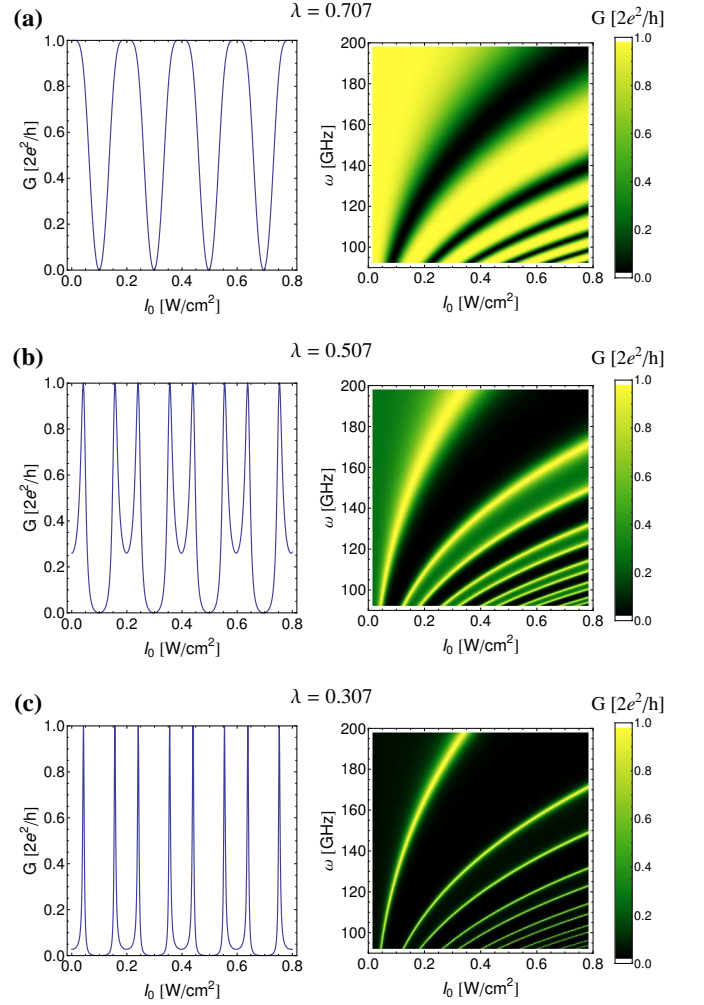


FIG. 2: (Color online) Conductance of a mesoscopic ring, G , under a circularly polarized electromagnetic wave as a function of wave intensity I_0 and wave frequency ω . Plots (a), (b) and (c) correspond to different transmission amplitudes λ between the current leads and the ring. Frames in the left column are fixed at the wave frequency $\omega = 100$ GHz. In all plots, the ring parameters are assumed to be $R = 10$ μm , $\varepsilon_F = 10$ meV, and $m_e = 0.1 m_{e0}$, where m_{e0} is the mass of free electron.

describing the AB-oscillations¹⁷ yields

$$G = \frac{2e^2}{h} \left[1 - \left| \frac{\sin^2(\Phi_{\text{eff}}/2\Phi_0)}{1 - \exp(i2\pi R k_F) \cos^2(\Phi_{\text{eff}}/2\Phi_0)} \right|^2 \right], \quad (23)$$

where k_F is the Fermi electron wave vector in the ring. For other amplitudes λ , the conductance G can be calculated numerically by using the same theory.^{15–17} Results of the calculations for different amplitudes λ are presented in Fig. 2.

For absolutely transparent QPCs ($\lambda = 0.707$), the regular AB-like oscillations take place (Fig. 2a). Decreasing the transparency (decreasing λ) changes the shape of the oscillation pattern (Figs. 2b and 2c). In the Fourier

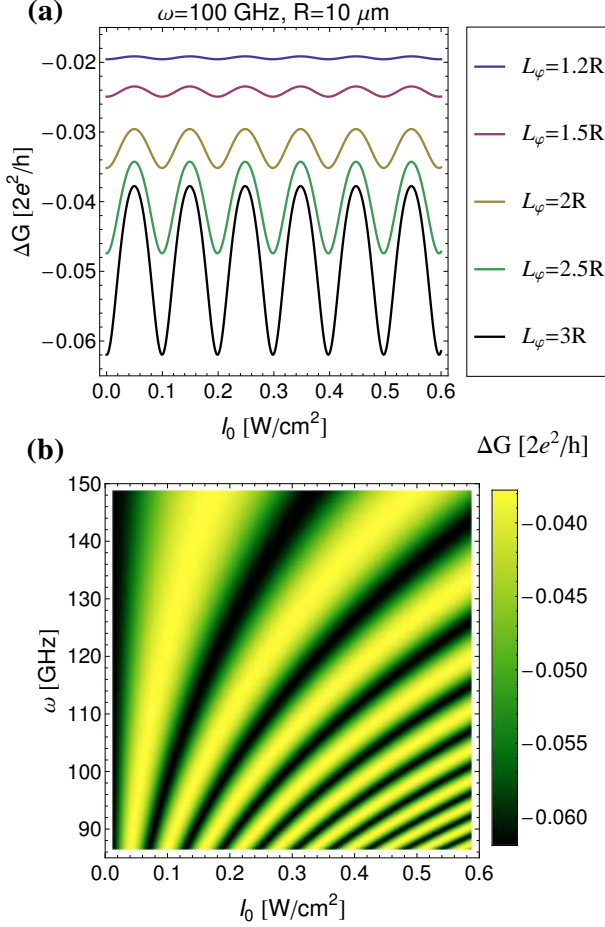


FIG. 3: (Color online) Weak-localization correction to the conductance of a mesoscopic ring, ΔG , under a circularly polarized electromagnetic field: (a) the correction is plotted as a function of field intensity I_0 for different values of L_φ with $\omega = 100$ GHz and $R = 10\mu\text{m}$; (b) the correction is plotted as a function of field intensity I_0 and of field frequency ω for $L_\varphi = 3R$ and $R = 10\mu\text{m}$.

spectrum of the conductance, the role of the higher harmonics increases, and eventually these harmonics with a half period become dominant (see Fig.2c). Physically, reduction of the period arises from an increased confinement of electrons inside the ring, caused by the decrease of transparency of the QPCs. This leads to an increase of the role of round trips of an electron inside the ring, which results in the increment of the effective electron path and, as a consequence, decrease of the period of the oscillations.

In the diffusive regime, the conductivity of a disordered ring-shaped conductor with the dephasing length L_φ can

be described by the expression

$$\Delta\sigma = -\frac{e^2 L_\varphi}{\pi^2 h} \frac{\sinh(2\pi R/L_\varphi)}{\cosh(2\pi R/L_\varphi) + \cos(4\pi\Phi_{\text{eff}}/\Phi_0)}, \quad (24)$$

which is derived from the conventional theory of AAS-oscillations⁷ by the replacement $\Phi \rightarrow \Phi_{\text{eff}}$.

The weak-localization correction to the conductance, $\Delta G = \Delta\sigma/\pi R$, is plotted in Fig. 3 for different values of the dephasing length L_φ . As expected, the correction oscillates with a period which is less than the period of AB-like oscillations (Fig. 2a) by a factor of 2. As for the amplitude of the oscillations, it decays exponentially when the dephasing length L_φ is much smaller than the distance between the QPCs, πR . Physically, this decay is caused by the electron waves losing their coherence quickly. It should be noted that an electromagnetic field can cause additional decoherence of electrons in conducting systems^{30–32} and, therefore, influences on the dephasing length L_φ . However, the condition of applicability of dressing field model, $\omega\tau \gg 1$, corresponds physically to the absence of energy exchange between conduction electrons and a dressing field, where τ is the characteristic electron relaxation time. Therefore, there is no heating of electrons by the field under this condition. As a consequence, the photon-induced breaking of phase coherence is negligibly small for a dressing field. Plotting the correction to the conductance, ΔG , in Fig. 3, we assumed the field to be high-frequency enough to neglect the phase decoherence arisen from the field.

Summarizing the aforesaid, we have shown that the interference of electron waves traveling through a mesoscopic ring exposed to a circularly polarized electromagnetic field is formally the same as in a ring subjected to a magnetic flux. As a consequence, the *optically-induced* Aharonov-Bohm effect appears. This effect manifests itself in the oscillating dependence of the ring conductance on the field intensity and field frequency. The periods of the optically-induced oscillations in the ballistic regime and the diffusive regime differ from each other by a factor of 2 in the same manner as periods of the oscillations induced by a magnetic flux. Therefore, the effect can be described formally in terms of the artificial $U(1)$ gauge field arisen from the strong electron-photon coupling.

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¹ Y. Aharonov and D. Bohm, Phys. Rev. **115**, 485 (1959).

- ² R. G. Chambers, Phys. Rev. Lett. **5**, 3 (1960).
- ³ R. A. Webb, S. Washburn, C. P. Umbach, and R. B. Laibowitz, Phys. Rev. Lett. **54**, 2696 (1985).
- ⁴ G. Timp, A. M. Chang, J. E. Cunningham, T. Y. Chang, P. Mankiewich, R. Behringer, and R. E. Howard, Phys. Rev. Lett. **58**, 2814 (1987).
- ⁵ P. G. N. de Vegvar, G. Timp, P. M. Mankiewich, R. Behringer, and J. Cunningham, Phys. Rev. B **40**, 3491 (1989).
- ⁶ B. J. van Wees, L. P. Kouwenhoven, C. J. P. M. Harmans, J. G. Williamson, C. E. Timmering, M. E. I. Broekaart, C. T. Foxon, and J. J. Harris, Phys. Rev. Lett. **62**, 2523 (1989).
- ⁷ B. L. Altshuler, A. G. Aronov, and B. Z. Spivak, Sov. Phys. JETP Lett. **33**, 94 (1981).
- ⁸ D. Y. Sharvin and Y. V. Sharvin, Sov. Phys. JETP Lett. **34**, 273 (1981).
- ⁹ G. Bergmann, Phys. Rev. B **28**, 2914 (1983).
- ¹⁰ B. Pannetier, J. Chaussy, R. Rammal, and P. Gandit, Phys. Rev. B **31**, 3209 (1985).
- ¹¹ O. V. Kibis, Phys. Lett. A **244**, 432 (1998).
- ¹² O. V. Kibis, Physica E (Amsterdam) **12**, 741 (2002).
- ¹³ O. V. Kibis, Physica E (Amsterdam) **13**, 699 (2002).
- ¹⁴ O. V. Kibis, Phys. Rev. Lett. **107**, 106802 (2011).
- ¹⁵ Y. Gefen, Y. Imry and M. Ya. Azbel, Phys. Rev. Lett. **52**, 129 (1984).
- ¹⁶ M. Büttiker, Y. Imry and M. Ya. Azbel, Phys. Rev. A **30**, 1982 (1984).
- ¹⁷ I. A. Shelykh, N. T. Bagraev, N. G. Galkin and L. E. Klyachkin, Phys. Rev. B **71**, 113311 (2005).
- ¹⁸ A. G. Aronov and Y. B. Lyanda-Geller, Phys. Rev. Lett. **70** 343 (1993).
- ¹⁹ A. F. Morpurgo, J. P. Heida, T. M. Klapwijk, B. J. van Wees, and G. Borghs, Phys. Rev. Lett. **80**, 1050 (1998).
- ²⁰ D. Frustaglia and K. Richter, Phys. Rev. B **69**, 235310 (2004).
- ²¹ T. Bergsten, T. Kobayashi, Y. Sekine, and J. Nitta, Phys. Rev. Lett. **97**, 196803 (2006).
- ²² O. V. Kibis, S. V. Malevanny, L. Hugget, D. G. W. Parfitt, M. E. Portnoi, Electromagnetics **25**, 425 (2005); O. V. Kibis, Phys. Lett. A **166**, 393 (1992).
- ²³ A. M. Alexeev, M. E. Portnoi, Phys. Rev. B **85**, 245419 (2012).
- ²⁴ V. B. Berestetskii, E. M. Lifshitz, L. P. Pitaevskii, *Quantum Electrodynamics* (Pergamon Press, Oxford, 1982).
- ²⁵ M.O. Scully and M.S. Zubairy, *Quantum Optics* (University Press, Cambridge, 2001).
- ²⁶ C. Cohen-Tannoudji, J. Dupont-Roc, G. Grynberg, *Atom-Photon Interactions: Basic Processes and Applications* (Wiley, Weinheim, 2004).
- ²⁷ O. V. Kibis, O. Kyriienko, I. A. Shelykh, Phys. Rev. B **87**, 245437 (2013).
- ²⁸ T. Arnold, C.-S. Tang, A. Manolescu, V. Gudmundsson, Phys. Rev. B **87**, 035314 (2013).
- ²⁹ J. Dalibard, F. Garbier, G. Juzeliūnas and P. Öhberg, Rev. Mod. Phys. **83**, 1523 (2011).
- ³⁰ B. L. Altshuler, A. G. Aronov, D. E. Khmel'nitsky, Solid State Commun. **39**, 619 (1981).
- ³¹ B. L. Altshuler, M. E. Gershenson, I. L. Aleiner, Physica E (Amsterdam) **3**, 58 (1998).
- ³² J.-S. Wenzler and P. Mohanty, Phys. Rev. B **77**, 121102(R) (2008).